Equation of State for Strange Quark Matter in a Separable Model

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Abstract

We present the thermodynamics of a nonlocal chiral quark model with separable 4-fermion interaction for the case of $U(3)$ flavor symmetry within a functional integral approach. The four free parameters of the model are fixed by the chiral condensate, and by the pseudoscalar meson properties (pion mass, kaon mass, pion decay constant). We discuss the $T = 0$ equation of state (EoS) which describes quark confinement (zero quark matter pressure) below the critical chemical potential $\mu_c = 333$ MeV. The new result of the present approach is that the strange quark deconfinement is separated from the light quark one and occurs only at a higher chemical potential of $\mu_{c,s} = 492$ MeV. We compare the resulting EoS to bag model ones for two and three quark flavors, which have the phase transition to the vacuum with zero pressure also at μ_c.

We study quark matter stars in general relativity theory assuming β-equilibrium with electrons and show that for configurations with masses close to the maximum of stability at $M = 1.62 \div 1.64 \ M_\odot$ strange quark matter can occur.

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I. INTRODUCTION

Since the discovery of the parton substructure of nucleons and its interpretation within the constituent quark model, much effort has been spent to explain the properties of these particles. The phenomenon of confinement, i.e. the property of quarks to exist only in bound states as mesons and baryons in all known systems, poses great difficulties for a describing theory. So far the problem has been solved by introducing a color interaction that binds all colored particles to "colorless" states.

However, it is believed and new experimental results [1] underline it, that at very high temperatures exceeding 150 MeV, or densities higher than three times nuclear matter density, a transition to deconfined quark matter can occur. Besides of the heavy ion collisions performed in particle physics, the existence of a deconfinement phase and its properties is of high importance for the understanding of compact stars [2] in astrophysics. These, that are popular as Neutron stars, imply core densities above three times the nuclear saturation density [4] so that quark matter is expected to occur in their interior [3] and several suggestions have been made in order to detect signals of the deconfinement transition [5,6].

Unfortunately, rigorous solutions of the fundamental theory of color interactions (Quantumchromodynamics (QCD)) for the EoS at finite baryon density could not be obtained yet, even Lattice gauge theory simulations have serious problems in this domain [7]. To describe interacting quark matter it is therefore necessary to find approximating models. The best studied one is the Nambu-Jona-Lasinio (NJL) model that was first developed to describe the interaction of nucleons [8] and has later been applied for modeling low-energy QCD [9–11] with particular emphasis on the dynamical breaking of chiral symmetry and the occurence of the pion as a quasi Goldstone boson. The application of the NJL model for studies of quark matter thermodynamics is problematic since it has no confinement and free quarks appear well below the chiral phase transition [12,13]. This contradicts to results from lattice gauge theory simulations of QCD thermodynamics where the critical temperatures for deconfinement and chiral restoration coincide. That can be helped by using a separable model which can be treated similarly to the NJL model but includes a momentum dependence for the interaction via formfactors. It has been shown [14] that in the chiral limit the model has no free quarks below the chiral transition.

Looking again at the densities of compact star cores and comparing with the results of NJL model calculations [13,15] it seems reasonable to include strange-flavor quarks in the model because the energy density is sufficiently high for their creation in weak processes. Therefore, we extend in the present work the separable model to the case of three quark flavors, assuming for simplicity $U(3)$ symmetry. We will calculate the partition function using the method of bosonisation and applying the mean-field approximation. Finally we will formulate the resulting thermodynamics of three-flavor quark matter and obtain numerical results for the quark matter EoS and compact star structure.

II. THE SEPARABLE QUARK MODEL

The starting point of our approach is an effective chiral quark model action with a four-fermion interaction in the current-current form

$$\mathcal{S}[q, \bar{q}] = \int \frac{d^4k}{(2\pi)^4} [\bar{q}(k) i(\not{k} + \hat{m}) q(k)$$

$$+ D_0 \int \frac{d^4k'}{(2\pi)^4} \sum_{\alpha=0}^{8} [j_s^\alpha(k) j_s^\alpha(k')$$

$$+ j_p^\alpha(k) j_p^\alpha(k')]] , \qquad (1)$$

where we restrict us here to the scalar current $j_s^\alpha(k) = \bar{q}(k) \lambda_\alpha f(k) q(k)$ and the pseudoscalar current $j_p^\alpha(k) = \bar{q}(k) i \gamma_5 \lambda_\alpha f(k) q(k)$ in Dirac space with $q(k)$ and $\bar{q}(k)$ being quark spinors and the formfactor $f(k)$ accounts for the nonlocality of the interaction. The action (1) is invariant under chiral rotations of the quark fields and color correlations are neglected (global color model). The generalization of previous models of this type [14,16,17] to the three-flavor case is done by using the $U(3)$ symmetry where λ_α are the Gell-Mann matrices and $\lambda_0 = \sqrt{\frac{2}{3}} I$.

Furthermore we do not include one of the possible models to account for the $U_A(1)$ anomaly since, at least in the quark representation of the Di Vecchia - Veneziano model [18], it can be shown that there is no contribution to the quark thermodynamics on the mean-field level [19]. For the quark mass matrix in flavor space we use the notation

$$\hat{m} = \sum_f m_f P_f , \qquad (2)$$

where m_f are the current quark masses and the projectors P_f on the flavor eigenstate $f = u, d, s$ are defined as

$$P_u = \frac{1}{\sqrt{6}} \lambda_0 + \frac{1}{2} \lambda_3 + \frac{1}{2\sqrt{3}} \lambda_8, \qquad (3)$$

$$P_d = \frac{1}{\sqrt{6}} \lambda_0 - \frac{1}{2} \lambda_3 + \frac{1}{2\sqrt{3}} \lambda_8, \qquad (4)$$

$$P_s = \frac{1}{\sqrt{6}} \lambda_0 - \frac{1}{\sqrt{3}} \lambda_8 \quad . \qquad (5)$$

Since the Matsubara frequencies in the $T \to 0$ limit become quasicontinuous variables, the summation over the fourth component k_4 of the 4-momentum has been replaced by the corresponding integration. According to the Matsubara formalism the calculations are performed in Euclidean space rather than in Minkowski space where we use $\gamma^4 = i\gamma^0$. The partition function in Feynman's path integral representation is given by

$$\mathcal{Z}[T, \hat{\mu}] = \int \mathcal{D}\bar{q} \mathcal{D}q \exp\left(\mathcal{S}[q, \bar{q}] - \int \frac{d^4k}{(2\pi)^4} i\hat{\mu}\gamma^4 \bar{q}q \right) , \qquad (6)$$

where the constraint of baryon number conservation is realized by the diagonal matrix of chemical potentials $\hat{\mu}$ (Lagrange multipliers) using the notation of the hat symbol analogous to (2).

In order to perform the functional integrations over the quark fields $\bar{q}$ and q we use the formalism of bosonisation (see [20] and references therein) which is based on the Hubbard-Stratonovich transformation of the four-fermion interaction terms employing the identity

$$\exp\left\{ D_0 \int_k \int_{k'} \sum_{\alpha=0}^{8} j_s^\alpha(k) j_s^\alpha(k') \right\} =$$

$$\mathcal{N} \prod_\alpha \int d\sigma^\alpha \exp\left[\frac{(\sigma^\alpha)^2}{4D_0} + \int_k \int_{k'} j_s^\alpha(k) \sigma^\alpha(k') \right] \ , \tag{7}$$

for the scalar and a similar one for the pseudoscalar channel, where for the phase space integral the abbreviation $\int_k = \int \frac{d^4k}{(2\pi)^4}$ has been used and $\mathcal{N}$ is a normalization factor. Now the generating functional is Gaussian in the quark fields and can be evaluated. We arrive at the transformed generating functional in terms of bosonic variables

$$\mathcal{Z}[T, \hat{\mu}] = \iint \mathcal{D}\sigma^\alpha \mathcal{D}\pi^\alpha \exp\left\{ \mathcal{S}[\sigma^\alpha, \pi^\alpha] \right\} \tag{8}$$

with the action functional

$$\mathcal{S}[\sigma^\alpha, \pi^\alpha] = -\int \frac{d^4k}{(2\pi)^4} \ln\left(\det_{DFC} [\tilde{k} + \hat{m} + \hat{\sigma} f(\tilde{k}) + i\gamma_5 \hat{\pi} f(\tilde{k})] \right)$$

$$+ \frac{\bar{\sigma}_\alpha \bar{\sigma}^\alpha}{4D_0} + \frac{\bar{\pi}_\alpha \bar{\pi}^\alpha}{4D_0} \ . \tag{9}$$

with analogous use of the already known hat symbol and the 4-vector $\tilde{k}_f = \left(\vec{k}, k_4 + i\mu_f \right)$. In order to further evaluate the integral over the auxiliary bosonic fields σ_α and π_α we expanded them around their mean values $\bar{\sigma}_\alpha$ and $\bar{\pi}_\alpha$ that minimize the action

$$\sigma_\alpha = \bar{\sigma}_\alpha + \tilde{\sigma}_\alpha(k)$$
$$\pi_\alpha = \bar{\pi}_\alpha + \tilde{\pi}_\alpha(k)$$

and neglect the fluctuations $\tilde{\sigma}_\alpha(k)$ and $\tilde{\pi}_\alpha(k)$ in the following. The mean values of the pseudoscalar field vanish for symmetry reasons [19]. The indices DFC refer to the determinant in Dirac-, flavor- and color- space. So we end up with the mean-field action

$$\mathcal{S}_{\text{MF}}[T, \{\mu_f\}] = \sum_f \left(-2N_c \int \frac{d^4k}{(2\pi)^4} [\ln(\tilde{k}_f^2 + M_f^2)] + \frac{\Delta_f^2}{8D_0} \right) \ , \tag{10}$$

with the effective quark masses $M_f = M_f(\tilde{k}) = m_f + \Delta_f f(\tilde{k})$ and the number of colors N_c. The flavor dependent mass gaps Δ_f are defined by $\hat{\sigma} = \sum_f \Delta_f P_f$.

A. Quark matter thermodynamics in mean field approximation

In the mean field approximation, the grand canonical thermodynamical potential is given by

$$\Omega(T, \{\mu\}) = \beta^{-1} \ln\left\{ \mathcal{Z}[T, \{\mu_f\}] / \mathcal{Z}[0, \{0\}] \right\}$$
$$= \beta^{-1} \left\{ \mathcal{S}_{\text{MF}}[T, \{\mu_f\}] - \mathcal{S}_{\text{MF}}[0, \{0\}] \right\} \ , \tag{11}$$

Where the divergent vacuum contribution has been subtracted. In what follows we consider the case $T = 0$ only. In order to interpret our result, we want to represent it as a sum of three terms

$$\Omega(0, \{\mu\}) = \sum_f \left(-2N_c \int \frac{d^4k}{(2\pi)^4} \ln \left(\frac{\tilde{k}_f^2 + M_f^2}{\tilde{k}_f^2 + m_f^2} \right) + \frac{\Delta_f^2}{8D_0} \right)$$

$$-2N_c \sum_f \left(\int \frac{d^4k}{(2\pi)^4} \ln \left(\frac{\tilde{k}_f^2 + m_f^2}{k_f^2 + m_f^2} \right) \right) \tag{12}$$

$$+ \sum_f \left(2N_c \int \frac{d^4k}{(2\pi)^4} \ln \left(\frac{k^2 + (M_f^0)^2}{k^2 + m_f^2} \right) - \frac{(\Delta_f^0)^2}{8D_0} \right) \ ,$$

where $M_f^0 = m_f + \Delta_f^0 f(k^2)$ are the effective quark masses in the vacuum. The second term on the r.h.s. of this equation is the renormalized thermodynamical potential of an ideal fermion gas [21]. The third term of Eq. (12) is independent of T and μ, i.e. it is a (thermodynamical) constant for the chosen model. Refering to the MIT bag model we call this term the bag-constant B. The remaining term includes the effects of quark interactions in the mean field approximation and can be evaluated numerically.

All thermodynamical quantities can now be derived from Eq. (12). For instance, pressure, density, energy density and the chiral condensate are given by:

$$pV = -\Omega \ , \ n = -\frac{\partial \Omega}{\partial \mu} \ , \ \varepsilon = -p + \mu n \ , \ <\bar{q}_f q_f> = \frac{\partial \Omega}{\partial m_q} \ . \tag{13}$$

Still the quark mass gaps Δ_f have to be determined. This is done by solving the gap equations which follow from the minimization conditions $\frac{\partial \Omega}{\partial \Delta_f} = 0$. The gap equations read

$$\Delta_f = 4D_0(-2N_c) \int \frac{d^4k}{(2\pi)^4} \frac{2M_f f(\tilde{k})}{\tilde{k}_f^2 + M_f^2} \ . \tag{14}$$

As can be seen from (14), for the chiral $U(3)$ quark model the three gap equations for Δ_u, Δ_d, Δ_s are decoupled and can be solved separately.

III. RESULTS FOR THE GAUSSIAN FORMFACTOR

A. Parametrization of the model

In the nonlocal separable quark model described above the formfactor of the interaction was not yet specified. In the following numerical investigations we will employ a simple Gaussian

$$f(k) = \exp(-k^2/\Lambda^2) \ , \tag{15}$$

which has been used previously for the description of meson [22] and baryon [23] properties in the vacuum as well as for those of deconfinement and mesons at finite temperature [17,24]. A systematic extension to other choices of formfactors can be found in [25,26].

The Gaussian model has five free parameters to be defined: the coupling constant D_0, the interaction range Λ, and the three current quark masses m_u, m_d, m_s. Setting $m_u = m_d =: m_q$ we restrict ourselves to four free parameters. These are fixed by the three well known

observables: pion mass $m_\pi = 140$ MeV, kaon mass $m_K = 494$ MeV and pion decay constant $f_\pi = 93$ MeV. The formulas for the meson masses and the decay constant are calculated as approximations of the Bethe-Salpeter equation including the generalized Goldberger-Treiman relation [26].

The fourth condition comes from values for the chiral condensate that are conform with phenomenology. The resulting parametrisations of the quark model are shown in Tab. I.

B. Thermodynamics for quark matter without β equilibrium

This case is relevant for systems which are considered for time scales larger than the typical strong interaction time of about 1 fm/c but smaller than the weak interaction time of several minutes, so that the presence of leptons (electrons) does not influence on the composition of quark matter and we can choose the chemical equilibrium with $\mu_u = \mu_d = \mu_s = \mu$. For the numerical calculations we choose the parameter set for the light quark condensate $-\langle \bar{u}u + \bar{d}d \rangle^{1/3} = 240$ MeV which is a typical value known from phenomenology. We consider the behavior of thermodynamical quantities at $T = 0$ with respect to the chemical potential. As we set $m_u = m_d$ earlier there is no difference between up- and down quarks and both are referred to as light quarks. Fig. 1 visualizes the behavior of the thermodynamical potential as a function of the light quark gap $\Delta_q = \Delta_u = \Delta_d$ for different values of the chemical potential μ. For $\mu < \mu_c = 333$ MeV the argument and the value of the global minimum is independent of μ which corresponds to a vanishing quark density (confinement). At the critical value $\mu = \mu_c = 333$ MeV a phase transition occurs from the massive, confining phase to a deconfining phase negligibly small mass gap. From the solution of the gap equation shown in Fig. 2 one can see that the strange quark gap of $\Delta_s = 682$ MeV still remains unchanged. Thus the strange quarks are confined until a higher value of the chemical potential $\mu_{c,s} = 492$ MeV is reached. This value is much bigger than the current strange quark mass. In Fig. 2 we separate by vertical lines the regions of full confinement, two-flavor deconfinement and full deconfinement. Thus, in the present model, the onset of strange quark deconfinement is inhibited. Moreover, the onset of a finite strange quark density is not determined by a drop in the strange gap which remains constant and even starts to rise for large μ values. This result of the present model drastically differs from those of bag models or NJL models. The reason is the 4-momentum dependence of the dynamical quark mass function which results in complex mass poles for the quark propagators and makes the naive identification of the mass gap with a real mass pole impossible [26].

The effect on thermodynamical quantities can be understood if we look at the pressure. In Fig. 3 we show for comparison the resulting equation of state for the pressure of the present separable model together with a two-flavor and a 3-flavor bag model. Both bag models are chosen such that the critical chemical potential for the deconfinement coincides with that of the separable model. The pressure of the present three-flavor separable model can be well described by a two-flavor bag model with a bag constant $B = 81.3$ MeV/fm^3 in the region of chemical potentials 333 MeV $\leq \mu \leq 492$ MeV where the third flavor is still confined. For comparison, the 3 flavor bag model has a bag constant $B = 100.7$ MeV/fm^3 and is considerably harder than the separable one due to the additional relatively light strange quark flavor.

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C. Inclusion of β equilibrium with electrons

Quark matter in β-equilibrium is to be supplemented with the two relations for conservation of baryon charge and electric charge. In the deconfined phase there are quarks and leptons (in our model case up, down, strange quarks and electrons) with vanishing net electric charge

$$Q_{\rm q}(\mu^u, \mu^d, \mu^s) + Q_{\rm L}(\mu^e) = \frac{2}{3}n_u - \frac{1}{3}(n_d + n_s) - n_e = 0 \ . \tag{16}$$

Taking into account the energy balance in weak interactions

$$d \leftrightarrow u + e^- + \bar{\nu}_e \tag{17}$$

$$s \leftrightarrow u + e^- + \bar{\nu}_e \tag{18}$$

and introducing the average quark chemical potential $\mu = \frac{1}{3}(\mu_u + \mu_d + \mu_s)$ we can write the β equilibrium conditions as

$$\mu_u = \mu - \frac{2}{3}\mu^e \ , \qquad \mu_d = \mu_s = \mu + \frac{1}{3}\mu^e \ . \tag{19}$$

Solving the equation of charge neutrality (16) one can find the chemical potential of electrons as a function of μ and using Eqs. (19) the equation of state can be given in terms of a single chemical potential μ. In Fig. 4 we show the composition of the three-flavor quark matter for the Gaussian separable model in the case of β equilibrium with electrons as a function of the energy density. As for the case without β equilibrium we can define also in Fig. 4 the regions of quark confinement ($\varepsilon < \varepsilon_c = 350$ MeV/fm^3) and three-flavor deconfinement ($\varepsilon > \varepsilon_{c,s} = 930$ MeV/fm^3). In the region of two-flavor deconfinement the concentrations of electrons, up- and down- quarks coincide with those of the two-flavor bag model except for the relatively small energies close to ε_c where the effect of a small dynamical quark mass leads to a density dependence of the composition. In Fig. 5 we demonstrate the influence of the β equilibrium on the equation of state. It can be seen that the difference between pressures with and without β equilibrium is limited to the region of intermediate densities, where the electron fraction reaches its maximum value $x_e \simeq 0.002$.

IV. APPLICATIONS FOR COMPACT STARS

One of the main goals for studying the strange quark matter equation of state is the possible application for compact stars. In particular, the hypothesis that strange quark matter might be more stable than ordinary nuclear matter [27] has lead to the investigation of possible consequences for properties of compact stars made thereof [28]. Most of these applications use the bag model equation of state where the result depends on the value of the bag constant as a free parameter. Recently, first steps have been made towards a description of strange quark matter within dynamical quark models such as the NJL model [29], where the parameters are fixed from hadron properties. The non-confining quark dynamics of this model, however, leads to predictions for dynamical quark masses and

critical parameters of the chiral phase transition which differ from those of confining models [30,14] and might be quantitatively incorrect. Here we want to extend previous studies of compact star properties with dynamically confining quark models to the strange quark sector and find the characteristics of stable compact star configurations with the equation of state derived above.

For the calculation of the self-bound configuration for the quark matter with gravitational interaction one needs the condition of mechanical equilibrium of the thermodynamical pressure with the gravitational force. This condition is given by the Tolman-Oppenheimer-Volkoff-Equation

$$\frac{dP}{dr} = -G(\varepsilon(r) + P(r))\frac{m(r) + 4\pi G r^3 P(r)}{r(r - 2Gm(r))} \tag{20}$$

and defines the profiles for all thermodynamical quantities in the case of nonrotating spherically symmetric distributed matter configurations in general relativity. In this equation m denotes the accumulated mass in the sphere with radius r given by

$$m(r) = 4\pi \int_0^r \varepsilon(r')r'^2 dr \tag{21}$$

The gravitational constant is denoted by G. The radius R of the star is defined by the condition that the pressure becomes zero on the surface of the star $P(R) = 0$. The total mass of the star is $M = m(R)$.

Each configuration has one independent parameter which could be chosen to be $\varepsilon(0)$, the central energy density. In Fig. 6 we show the dependence of the total mass of the configuration as a function of the central density and the radius for the separable quark model and for the bag model in the cases of two and three flavors respectively.

The rising branches of the mass-radius or mass-density relations correspond to the families of stable compact stars. The maximum possible mass for the separable model is 1.64 $M_\odot$ for the three-flavor and 1.71 $M_\odot$ for the two-flavor case. The maximal central density is about 1350 MeV/fm^3 which allows for the three-flavor case to have strange quark matter in the core of the quark star. The comparison with the corresponding bag model strange stars shows that the latter are more compact, their maximum radius is about 8 km, and less massive with a maximum mass of about 1.5 $M_\odot$. The maximum radius of stars within the separable model is 11 km and thus exceeds the radii for both two and three flavor bag model quark stars. The origin of this difference is the behavior of the pressure in the low density region.

V. CONCLUSION

For neutron stars it is relevant to include the effects of strange flavor in a model for quark matter. In the simplest case considering $U(3)$ symmetry this can be done without increasing the complexity of the generating functional. We showed that in our separable model the gap equations decouple and can be solved separately. The resulting thermodynamics can be solved numerically and gives the equations of state for interacting quark matter. Unlike the well known NJL model the separable model is able to express the effects of confinement

in the thermodynamical quantities. The new result obtained within the present approach is the separation of the deconfinement of light quark flavors at $\mu_c = 333$ MeV from that of strange quarks which occurs only at a higher chemical potential of $\mu_{c,s} = 492$ MeV. A consequence for the application of the EoS presented here in compact star calculations is that strange quarks do occur only close to the maximum mass of 1.64 $M_\odot$, i.e. that for masses below 1.62 $M_\odot$, only two-flavor quark matter can occur.

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TABLES

$-<\bar{u}u + \bar{d}d>^{1/3}$ [MeV]	Λ [MeV]	D_0 [GeV^{-2}]	m_q [MeV]	m_s [MeV]	Δ_q^0 [MeV]	Δ_s^0 [MeV]
230	659.2	29.32	6.8	143.5	549.9	767.8
235	697.6	23.88	6.4	136.1	497.0	719.3
240	736.5	19.88	6.0	129.5	453.8	682.1
245	775.5	16.85	5.6	123.4	419.7	653.1
250	814.9	14.49	5.3	118.0	391.1	630.0
255	853.8	12.66	5.0	112.9	368.7	611.4
260	894.2	11.11	4.7	108.1	349.1	596.1

TABLE I. Parameter sets for the Gaussian separable model for different values of the chiral condensate $<\bar{u}u + \bar{d}d>$.

FIG. 1. Dependence of the thermodynamical potential on the light flavor gap $\Delta_q = \Delta_u = \Delta_d$ (order parameter) for different values of the chemical potential, $\Delta_s = 682$ MeV.
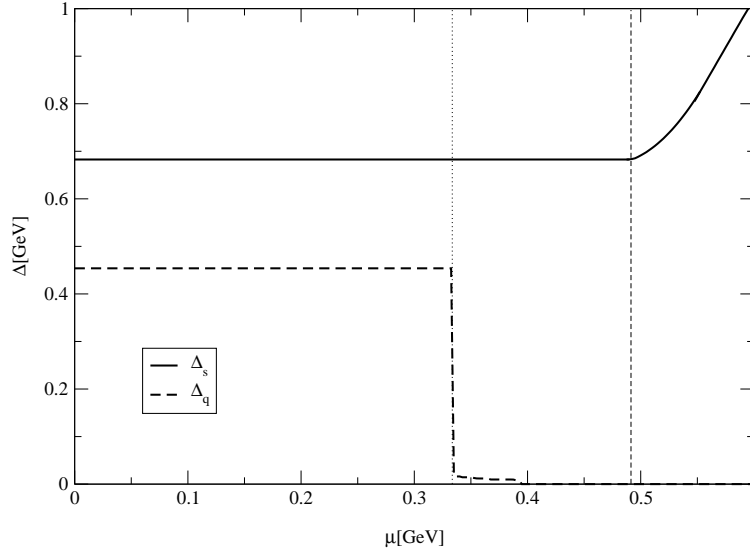
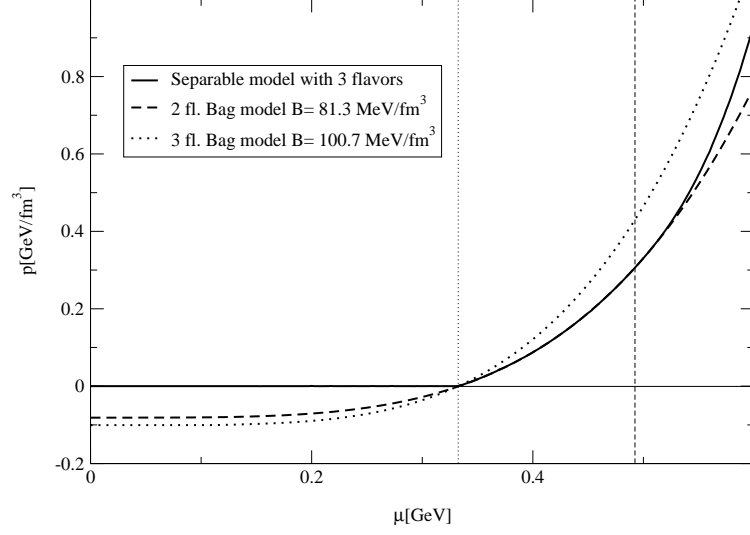FIG. 2. Solutions of the gap equations that minimize the potential

FIG. 3. Pressure of the quark matter as a function of the chemical potential for the separable model (solid line) compared to a three-flavor (dotted line) and a two-flavor (dashed line) bag model. All models have the same critical chemical potential $\mu_c = 333$ MeV for (light) quark deconfinement.
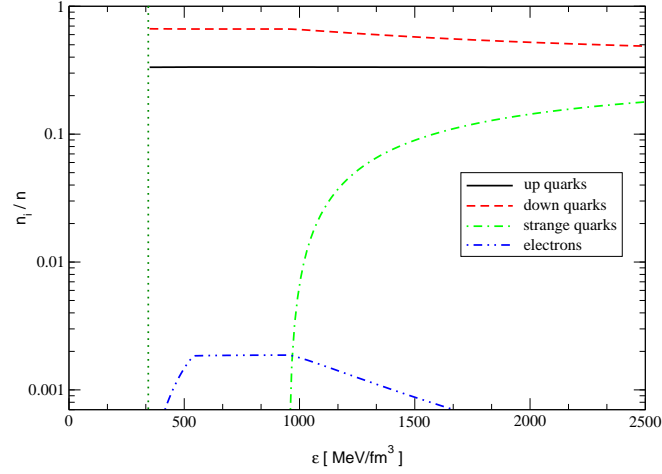

FIG. 4. Composition of three-flavor quark matter in β equilibrium with electrons.
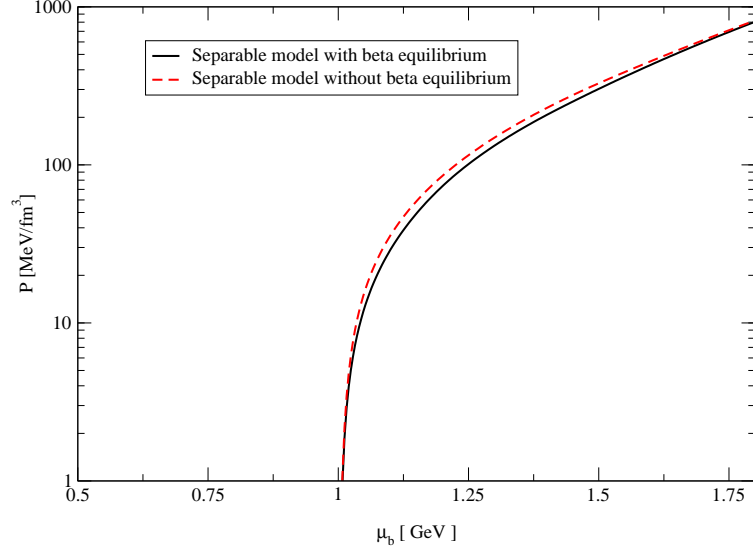
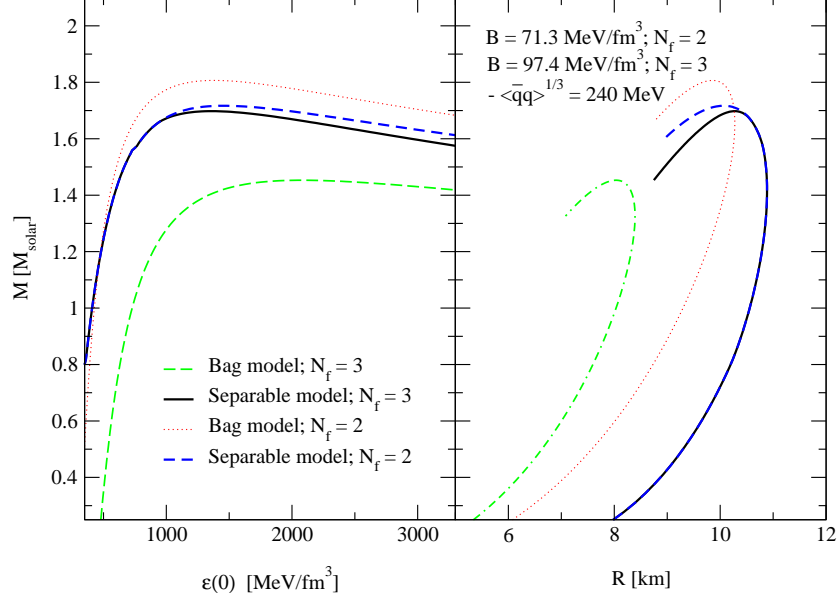FIG. 5. Pressure of three-flavor quark matter with β equilibrium and without.

FIG. 6. Stability for compact stars composed of quark matter.